

Sneutrino-Antisneutrino Mixing and Neutrino Mass in Anomaly-mediated Supersymmetry Breaking Scenario

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Abstract

In supersymmetric models with nonzero Majorana neutrino mass, the sneutrino and antisneutrino mix, which may lead to same sign dilepton signals in future collider experiments. We point out that the anomaly-mediated supersymmetry breaking scenario has a good potential to provide an observable rate of such signals for the neutrino masses suggested by the atmospheric and solar neutrino oscillations. The sneutrino mixing rate is naturally enhanced by $m_{3/2}/m_{\tilde{\nu}} = \mathcal{O}(4\pi/\alpha)$ while the sneutrino decay rate is small enough on a sizable portion of the parameter space. We point out also that the sneutrino-antisneutrino mixing can provide much stronger information on some combinations of the neutrino masses and mixing angles than neutrino experiments.

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Current data from the atmospheric and solar neutrino experiments strongly suggest that the neutrinos have small but nonzero masses [1]. As was pointed out in Refs. [2–5], in supersymmetric models with nonzero Majorana neutrino mass, the sneutrino ($\tilde{\nu}$) and antineutrino ($\tilde{\nu}^*$) mix to each other. The mixing rate is generically given by

$$\Delta m_{\tilde{\nu}} = C_{\nu} m_{\nu} / m_{\tilde{\nu}} \quad (1)$$

where m_{ν} is the lepton-number violating ($\Delta L = 2$) Majorana neutrino mass, $m_{\tilde{\nu}}$ is the lepton-number conserving ($\Delta L = 0$) sneutrino mass, and C_{ν} is determined mainly by the soft supersymmetry (SUSY) breaking in the lepton-number violating sector of the underlying theory. The atmospheric and solar neutrino data give the neutrino square mass differences $\Delta m_{\nu}^2 \lesssim \mathcal{O}(10^{-3}) \text{ eV}^2$, so it is quite unlikely that any of the neutrino masses is significantly bigger than $\mathcal{O}(1) \text{ eV}$ even when we include the possibility of nearly degenerate neutrino masses. Also the consideration of radiative corrections to m_{ν} induced by $\Delta m_{\tilde{\nu}}$ [3] leads to the bound $\Delta m_{\tilde{\nu}} / m_{\nu} \lesssim \mathcal{O}(4\pi/\alpha)$, implying $\Delta m_{\tilde{\nu}} \lesssim \mathcal{O}(1) \text{ keV}$. Such a small mixing rate can be probed by the $\tilde{\nu}$ - $\tilde{\nu}^*$ oscillation which would result in same sign dilepton signals when the sneutrino pairs decay into charged leptons.

To have an observable rate of same sign dilepton signals, $\tilde{\nu}$ must have an enough time to mix with $\tilde{\nu}^*$ before it decays. For this, we need the sneutrino decay width $\Gamma_{\tilde{\nu}} \lesssim \mathcal{O}(1) \text{ keV}$ in view of $\Delta m_{\tilde{\nu}} \lesssim \mathcal{O}(1) \text{ keV}$. Such a small decay rate would not be possible if the two body decay channel $\tilde{\nu} \rightarrow \nu \tilde{\chi}^0$ or $\ell^- \tilde{\chi}^+$ is open for the neutralinos $\tilde{\chi}^0$ or charginos $\tilde{\chi}^+$. It was pointed out in [3] that the most plausible scenario for $\Gamma_{\tilde{\nu}} \lesssim \mathcal{O}(1) \text{ keV}$ is to have

$$m_{\tilde{\tau}_1} < m_{\tilde{\nu}} < m_{\tilde{\chi}^0}, m_{\tilde{\chi}^+}, \quad (2)$$

where $\tilde{\tau}_1$ denotes the lighter stau. Then sneutrinos decay mainly into three-body final states with sizable branching ratio into a charged lepton: $\tilde{\nu} \rightarrow \ell^- \tilde{\tau}_1^+ \nu_{\tau}, \nu \tilde{\tau}_1^{\pm} \tau^{\mp}$ with $\Gamma_{\tilde{\nu}} \lesssim \mathcal{O}(1) \text{ keV}$. The mass hierarchy (2) would mean that the stau is the lightest supersymmetric particle in the minimal supersymmetric standard model (MSSM) sector, which would not be cosmologically allowed if it is stable. This difficulty can be easily avoided if one assumes

a light singlet fermion ψ which has very weak couplings to the MSSM sector, e.g. a light gravitino or axino, with which $\tilde{\tau}_1$ decays into $\tau\psi$. Alternatively, one may introduce a tiny R -parity violating coupling which would trigger $\tilde{\tau}_1 \rightarrow \ell\nu$. Still $\tilde{\tau}_1$ can live long enough inside the detector, so clearly distinguished from other charged sleptons. Then the same sign dilepton events induced by the $\tilde{\nu}\text{-}\tilde{\nu}^*$ mixing are accompanied by long-lived same sign stau pairs, so provide a rather clean signal for the $\tilde{\nu}\text{-}\tilde{\nu}^*$ mixing if we focus on the final states involving $\ell\ell\tilde{\tau}_1\tilde{\tau}_1$ for $\ell = e, \mu$.

Obviously, for a given neutrino mass, models with bigger $C_\nu/m_{\tilde{\nu}} = \Delta m_{\tilde{\nu}}/m_\nu$ have better prospect for observable $\tilde{\nu}\text{-}\tilde{\nu}^*$ mixing. In this paper, we point out that the anomaly-mediated SUSY breaking (AMSB) scenario [6] generically predicts $C_\nu/m_{\tilde{\nu}} = \mathcal{O}(4\pi/\alpha) \gg 1$ if the neutrino mass is generated by SUSY preserving dynamics at high energy scale as in the conventional seesaw model [7]. Furthermore the mass hierarchy (2) can be obtained on a sizable portion of the phenomenologically allowed parameter space of the AMSB model. These features are not shared by the minimal supergravity (SUGRA) model [8] or the gauge-mediated SUSY breaking (GMSB) model [9], so the AMSB model has much better potential to provide observable same sign dilepton signals induced by the $\tilde{\nu}\text{-}\tilde{\nu}^*$ mixing than other SUSY breaking models. An interesting feature of the $\tilde{\nu}\text{-}\tilde{\nu}^*$ mixing in the AMSB scenario is that it provides rather strong information on the neutrino mass matrix elements $(m_\nu)_{ii} = \sum_a U_{ia}^2 m_{\nu_a}$ where $i = e, \mu, \tau$ denote the flavor eigenstates while $a = 1, 2, 3$ stand for the neutrino mass eigenstates with the MNS mixing matrix U_{ia} . Atmospheric neutrino data suggests $(m_\nu)_{\mu\mu} \simeq 3 \times 10^{-2}$ eV for hierarchical neutrino masses, while $(m_\nu)_{\mu\mu}$ can be bigger if the neutrino masses are approximately degenerate. Then the same sign dimuon events from $\tilde{\nu}_\mu\text{-}\tilde{\nu}_\mu^*$ mixing can be used to distinguish $(m_\nu)_{\mu\mu} = 3 \times 10^{-2}$ eV from a bigger value of $(m_\nu)_{\mu\mu}$. Also the same sign dielectron events from $\tilde{\nu}_e\text{-}\tilde{\nu}_e^*$ mixing can be used to probe $(m_\nu)_{ee}$ down to the order of 10^{-4} eV which is much smaller than the current bound on $(m_\nu)_{ee}$ from ν -less double beta decays. In the following, we will examine these points in more detail.

Let us first discuss some generic features of the AMSB model. Anomaly mediation assumes that supersymmetry breaking in the hidden sector is transmitted to the MSSM

fields *mainly* by the Weyl compensator superfield Φ_0 of the supergravity multiplet [6]:

$$\Phi_0 = 1 + \theta^2 M_{\text{aux}}, \quad (3)$$

where M_{aux} is generically of order the gravitino mass $m_{3/2}$. The couplings of Φ_0 to generic matter multiplets are determined by the super-Weyl invariance. Therefore at classical level, Φ_0 is coupled to the MSSM fields only through *dimensionful* supersymmetric parameters, while the Φ_0 -couplings through dimensionless parameters arise from radiative corrections. More explicitly, a super-Weyl invariant effective action can be written as

$$S_{\text{eff}} = \int d^4x d^4\theta \left[Z_I(Q/\sqrt{\Phi_0\Phi_0^*})\Phi_I^*\Phi_I + \frac{1}{8}g_a^{-2}(Q/\sqrt{\Phi_0\Phi_0^*})DV^a\bar{D}^2DV^a \right] \\ + \left[\int d^4x d^2\theta \left(y_{IJK}\Phi_I\Phi_J\Phi_K + \frac{\Phi_0}{M}\gamma_{IJKL}\Phi_I\Phi_J\Phi_K\Phi_L \right) + \text{h.c.} \right] \quad (4)$$

where Q denotes the renormalization scale, D and \bar{D} are the supercovariant derivatives on the real gauge superfields V_a , and Φ_I are the chiral matter superfields. Here the quartic terms in the superpotential are assumed to be induced by supersymmetry preserving dynamics, e.g. by the exchange of heavy particles with supersymmetric mass M . Note that such heavy particles can be integrated out *while preserving the super-Weyl invariance*.

With Φ_0 given as (3), S_{eff} would determine the pure anomaly-mediated soft parameters in a manner which is valid at an arbitrary scale Q . However it predicts a tachyonic slepton, so one needs some additional source of SUSY breaking. The simplest possibility is an additional universal soft scalar square mass m_0^2 introduced at the GUT scale $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV. This defines the minimal AMSB model which predicts the following forms of soft supersymmetry breaking terms [6,10]:

$$\mathcal{L}_{\text{soft}} = m_I^2|\phi_I|^2 + \left(\frac{1}{2}M_a\lambda^a\lambda^a + A_{IJK}y_{IJK}\phi_I\phi_J\phi_K + \frac{C_{IJKL}\gamma_{IJKL}}{M}\phi_I\phi_J\phi_K\phi_L + \text{h.c.} \right) \quad (5)$$

where the gaugino masses M_a , the soft scalar masses m_I and the soft coefficients A_{IJK}, C_{IJKL} are given by

$$M_a = -\frac{b_a\alpha_a}{4\pi}M_{\text{aux}}, \quad m_I^2 = -\frac{1}{4}\frac{d\gamma_I}{d\ln Q}|M_{\text{aux}}|^2 + m_0^2,$$

$$\begin{aligned}
A_{IJK} &= \frac{1}{2}(\gamma_I + \gamma_J + \gamma_K)M_{\text{aux}}, \\
C_{IJKL} &= \frac{1}{2}(2 + \gamma_I + \gamma_J + \gamma_K + \gamma_L)M_{\text{aux}}.
\end{aligned} \tag{6}$$

Here $b_a = (3, -1, -33/5)$ are the one-loop beta function coefficients for $SU(3)_c \times SU(2)_L \times U(1)_Y$ in the GUT normalization and $\gamma_I = d \ln Z_I / d \ln Q$ are the anomalous dimension of Φ_I . Note that still the expressions of M_a , A_{IJK} and C_{IJKL} are valid at arbitrary scale, while the expression of m_I^2 is valid *only* at M_{GUT} .

Applying the above results to the sneutrino-antisneutrino mixing is rather straightforward. To be specific, we will assume that the neutrino masses are generated (mainly) by supersymmetry preserving dynamics at an energy scale M far above the weak scale. This high energy dynamics may be the exchange of heavy singlet neutrino with mass M [7], or the exchange of heavy triplet Higgs boson [11], or some stringy dynamics. Independently of its detailed shape, this high energy dynamics can be integrated out while preserving the super-Weyl invariance. Then at the weak scale, the theory can be described by an effective superpotential including the super-Weyl invariant dimension 5 operators for neutrino masses and also the associated soft SUSY breaking terms,

$$\begin{aligned}
\Delta W_{\text{eff}} &= \frac{\Phi_0}{M} \gamma_{ij} (L_i H_2)(L_j H_2), \\
\Delta \mathcal{L}_{\text{soft}} &= \frac{C_{ij} \gamma_{ij}}{M} (\tilde{\ell}_i h_2)(\tilde{\ell}_j h_2),
\end{aligned} \tag{7}$$

where L_i ($i = e, \mu, \tau$) and H_α ($\alpha = 1, 2$) denote the lepton and Higgs doublet superfields with the scalar components $\tilde{\ell}_i$ and h_α , respectively, and $C_{ij} \simeq M_{\text{aux}}$.

After the electroweak symmetry breaking, ΔW_{eff} gives a neutrino mass matrix $(m_\nu)_{ij} = 2\langle h_2 \rangle^2 \gamma_{ij} / M$. Including the contribution from $\Delta \mathcal{L}_{\text{soft}}$, the sneutrino masses are given by

$$(m_\nu^2)_{ij} \tilde{\nu}_i^* \tilde{\nu}_j + \left\{ \frac{1}{2} (\Delta m_\nu^2)_{ij} \tilde{\nu}_i \tilde{\nu}_j + \text{h.c.} \right\}, \tag{8}$$

where the sneutrino square mass matrix $(m_\nu^2)_{ij} \simeq m_\nu^2 \delta_{ij}$ with $m_\nu^2 = \frac{1}{2} M_Z^2 \cos 2\beta + m_\ell^2$ for the slepton doublet square mass matrix $(m_\ell^2)_{ij} \simeq m_\ell^2 \delta_{ij}$ and the $\Delta L = 2$ sneutrino square mass matrix is given by $(\Delta m_\nu^2)_{ij} = (C_{ij} + 2\mu \cot \beta)(m_\nu)_{ij}$. The $\tilde{\nu}$ - $\tilde{\nu}^*$ mixing rate is determined

by the sneutrino mass-splitting $\Delta m_{\tilde{\nu}} = \Delta m_{\tilde{\nu}}^2/m_{\tilde{\nu}}$. In this regard, a *distinctive feature* of the AMSB model is that $C_{ij} \simeq M_{\text{aux}}$ is induced at tree level while $m_{\tilde{\nu}}$ is loop-suppressed, so $\Delta m_{\tilde{\nu}}$ is enhanced (relative to m_{ν}) by the factor $M_{\text{aux}}/m_{\tilde{\nu}} = \mathcal{O}(4\pi/\alpha)$:

$$(\Delta m_{\tilde{\nu}})_{ij} \simeq (m_{\nu})_{ij} M_{\text{aux}}/m_{\tilde{\nu}} = \mathcal{O}(4\pi m_{\nu}/\alpha). \quad (9)$$

Consequently, for a given neutrino mass, the AMSB model has better potential to give a sizable $\tilde{\nu}$ - $\tilde{\nu}^*$ mixing than other models with $C_{ij}/m_{\tilde{\nu}} = \mathcal{O}(1)$. Furthermore, as can be inferred from Fig. 1, a significant portion of the phenomenologically viable parameter space of the minimal AMSB model leads to the mass hierarchy (2), which is a feature *not* shared by the minimal SUGRA or GMSB models. This is partly because in the AMSB model the lightest neutralino ($\tilde{\chi}_1^0$) and the lightest chargino ($\tilde{\chi}_1^\pm$) are nearly degenerate and the mass gap between the sleptons and $\tilde{\chi}_1^0, \tilde{\chi}_1^\pm$ is narrower than the minimal SUGRA and GMSB models [10].

Taking an analogy to the B -meson mixing, it is straightforward to compute the probability for a $\tilde{\nu}$ - $\tilde{\nu}^*$ pair produced in e^+e^- collider to yield same-sign dilepton signal [3,12]. The amplitude for $e^+(p_1) + e^-(p_2) \rightarrow \tilde{\nu}_i(q_1) + \tilde{\nu}_i^*(q_2)$ is easily computed to be

$$A_i = \frac{1}{2} g^2 \bar{v}(p_2) (\not{p}_1 - \not{p}_2) (X_i P_L + Y_i P_R) u(p_1), \quad (10)$$

where $P_{L,R} = (1 \pm \gamma_5)/2$, $X_i = K_Z(s_W^2 - \frac{1}{2})/c_W^2 + \delta_{ie} \sum_n |V_{n1}|^2 K_n$, $Y_i = K_Z s_W^2/c_W^2$ for the chargino ($\tilde{\chi}_n^\pm$) mixing matrix V_{nm} , $K_Z = 1/(s - M_Z^2)$, $K_n = 1/[(p_1 - q_1)^2 - m_{\chi_n}^2]$, $c_W = \cos \theta_W$, $s_W = \sin \theta_W$, and the c.m. energy \sqrt{s} . Here K_Z represents the contributions from the Z boson mediated s -channel diagrams, while K_n is from the chargino mediated t -channel diagrams. With this amplitude, the initial $\tilde{\nu}$ - $\tilde{\nu}^*$ state is given by

$$|\tilde{\nu}\tilde{\nu}^*; 0\rangle = \sum_i \alpha_i |\tilde{\nu}_i(\vec{q})\rangle |\tilde{\nu}_i^*(-\vec{q})\rangle + \beta_i |\tilde{\nu}_i^*(\vec{q})\rangle |\tilde{\nu}_i(-\vec{q})\rangle, \quad (11)$$

where $\alpha_i = A_i(\vec{q}_1 = \vec{q})$, $\beta_i = A_i(\vec{q}_2 = \vec{q})$ (up to overall normalization) and the momentum vector \vec{q} spans only the upper hemisphere, i.e. $\cos \theta \geq 0$ for the angle θ between e^- and $\tilde{\nu}$ flight directions.

With (2), the sneutrinos decay as $\tilde{\nu} \rightarrow \ell^- \tilde{\tau}_1^+ \nu_\tau, \nu \tilde{\tau}_1^\pm \tau^\mp$. It turns out that, in most of the parameter space yielding an observable rate of same sign dilepton signals, these decays are induced dominantly by the chargino or neutralino exchange, so the decay widths are (approximately) flavor-independent. Then the effective Hamiltonian determining the evolution (11) can be written as

$$H_{\text{eff}} = \begin{pmatrix} (m_{\tilde{\nu}} - \frac{i}{2}\Gamma_{\tilde{\nu}})\delta_{ij} + (\delta m_{\tilde{\nu}})_{ij} & \frac{1}{2}(\Delta m_{\tilde{\nu}})_{ij} \\ \frac{1}{2}(\Delta m_{\tilde{\nu}})_{ij}^* & (m_{\tilde{\nu}} - \frac{i}{2}\Gamma_{\tilde{\nu}})\delta_{ij} + (\delta m_{\tilde{\nu}})_{ij}^* \end{pmatrix}, \quad (12)$$

where $\delta m_{\tilde{\nu}}$ represents the deviation from the exact degeneracy of the $\Delta L = 0$ sneutrino masses. Since $\delta m_{\tilde{\nu}} \gg \Delta m_{\tilde{\nu}}$ in our case, it is most convenient to describe the $\tilde{\nu}$ - $\tilde{\nu}^*$ mixing in the field basis in which $\delta m_{\tilde{\nu}}$ is diagonal. In the AMSB scenario, the charged lepton mass matrix can be diagonalized simultaneously with $\delta m_{\tilde{\nu}}$. In such field basis, we find the probability P_i for the initial state (11) to produce a same sign dilepton $\ell_i^- \ell_i^-$ or $\ell_i^+ \ell_i^+$:

$$P_i = \frac{1}{\sum_i \sigma_i} \int d\Phi_2 \frac{1}{8s} \frac{B_i^2}{(1+x_i^2)^2} \left\{ \frac{1}{2}(|\alpha_i|^2 + |\beta_i|^2)(2+x_i^2)x_i^2 + \text{Re}(\alpha_i^* \beta_i)x_i^2 \right\}, \quad (13)$$

where $B_i = \text{Br}(\tilde{\nu}_i \rightarrow \ell_i X)$, $x_i = |(\Delta m_{\tilde{\nu}})_{ii}|/\Gamma_{\tilde{\nu}} = |(m_\nu)_{ii}|M_{\text{aux}}/m_{\tilde{\nu}}\Gamma_{\tilde{\nu}}$ and σ_i denotes the total cross section for $e^+e^- \rightarrow \tilde{\nu}_i \tilde{\nu}_i^*$. Here the 2-body phase space integration ($d\Phi_2$) for the initial $\tilde{\nu}\tilde{\nu}^*$ is performed for $\cos\theta \geq 0$ and we have ignored the effects suppressed by $\Delta m_{\tilde{\nu}}/\delta m_{\tilde{\nu}}$.

The same sign dilepton probability (13) shows that in the AMSB scenario the $\tilde{\nu}$ - $\tilde{\nu}^*$ mixing provides information on the neutrino matrix elements $(m_\nu)_{ii} = \sum_a U_{ia}^2 m_{\nu_a}$ where m_{ν_a} and U_{ia} denote the neutrino mass eigenvalues and the MNS mixing matrix, respectively. This would be true in other SUSY breaking models as long as the SUSY breaking is transmitted to the observable sector in a *flavor-blind* way. Currently $(m_\nu)_{ee}$ is bounded to be less than 0.2 eV by the ν -less double beta decay and this bound can be relaxed by a factor of few due to the uncertainty in the involved nuclear matrix elements. As we will see, the $\tilde{\nu}$ - $\tilde{\nu}^*$ mixing allows us to probe $(m_\nu)_{ee}$ down to the order of 10^{-4} eV in the AMSB scenario. Various neutrino oscillation experiments including the atmospheric and solar neutrino oscillations provide information on $m_{\nu_a}^2 - m_{\nu_b}^2$ and U_{ia} , for instance $\Delta m_{\text{atm}}^2 = |m_{\nu_3}^2 - m_{\nu_2}^2|^2 \simeq 3 \times 10^{-3} \text{ eV}^2$ and $|U_{\mu 3}| \simeq |U_{\tau 3}| = 1/\sqrt{2}$. Still the information on $(m_\nu)_{ii}$ from the $\tilde{\nu}$ - $\tilde{\nu}^*$ mixing are different

from the information on neutrino masses and mixing angles from neutrino oscillations. So the $\tilde{\nu}$ - $\tilde{\nu}^*$ mixing can provide information on neutrino masses and mixing angles which are complementary to those from the neutrino experiments.

Same sign dilepton events may come also from the pair-produced neutralinos which would decay as $\tilde{\chi}^0 \rightarrow \ell\tilde{\ell}$. However as long as we focus on the e and μ flavors the same sign dileptons from the neutralino pair accompany \tilde{e} or $\tilde{\mu}$, so can be clearly distinguished from those from the $\tilde{\nu}$ - $\tilde{\nu}^*$ mixing accompanying the $\tilde{\tau}_1$ pair. With this observation, we performed a numerical analysis to find the parameter region of the minimal AMSB model yielding a sizable number of same sign dilepton events per year, N_i ($i = e, \mu$), for a future e^+e^- linear collider with the integrated luminosity 500 fb^{-1} at $\sqrt{s} = 500 \text{ GeV}$. As usual, we replace the Higgs μ and B parameters by M_Z and $\tan\beta$ under the condition of electroweak symmetry breaking. Using the standard RG analysis, the superparticle mass spectrums are obtained to compute $\Gamma_{\tilde{\nu}}$ on the parameter region of $(m_0, M_{\text{aux}}, \tan\beta)$ leading to the mass hierarchy (2). When $\tan\beta$ increases for a given M_{aux} , the mass hierarchy (2) requires a larger m_0 leading to a larger $m_{\tilde{\nu}}$. On the other hand, the enhanced left-right mixing gives an effect to lower $m_{\tilde{\tau}_1}$, so the net result is to increase $m_{\tilde{\nu}}/m_{\tilde{\tau}_1}$. The phase space of the three body decays $\tilde{\nu} \rightarrow \ell\tilde{\tau}_1\nu_\tau, \nu\tilde{\tau}_1\tau$ is highly sensitive to $m_{\tilde{\nu}}/m_{\tilde{\tau}_1}$. As a result, for a given neutrino mass, $\Gamma_{\tilde{\nu}}$ is a sharply increasing function of $\tan\beta$, so small $\tan\beta$ is favored for sizable N_i . From a detailed numerical analysis, we find that $\tan\beta \lesssim 10$ is required to have a sizable N_i for $m_\nu \lesssim \mathcal{O}(1) \text{ eV}$.

About the values of $(m_\nu)_{\mu\mu}$, we considered two cases. In the first, neutrino masses are assumed to be hierarchical, which would give $(m_\nu)_{\mu\mu} \simeq U_{\mu 3}^2 \sqrt{\Delta m_{\text{atm}}^2} \simeq 3 \times 10^{-2} \text{ eV}$, while in the second case neutrino masses are assumed to be approximately degenerate with $(m_\nu)_{\mu\mu} = 0.3 \text{ eV}$. We find that the number of same sign dimuon events per year (N_μ) for the first case is bigger than the value for the second case by about factor 5, so hierarchical and (approximately) degenerate neutrino masses are clearly distinguished from each other. In Fig. 1, we depict the parameter regions with $\tan\beta = 5$ yielding $N_\mu \geq 20, 10^2, 5 \times 10^2$ for $(m_\nu)_{\mu\mu} = 3 \times 10^{-2} \text{ eV}$ and $N_\mu \geq 10^2, 5 \times 10^2, 2 \times 10^3$ for $(m_\nu)_{\mu\mu} = 0.3 \text{ eV}$. We also

searched for the parameter regions with $\tan\beta = 5$ yielding $N_e \geq 20, 10^2, 5 \times 10^2$ for $(m_\nu)_{ee} = 10^{-2}, 10^{-3}, 10^{-4}$ eV and depict the results in Fig. 2. Note that the t -channel contribution to $e^+e^- \rightarrow \tilde{\nu}_e\tilde{\nu}_e^*$ enhances N_e relative to N_μ , so that we can have a sizable N_e even for $(m_{ee}) = 10^{-4}$ eV.

As available constraints on the model, we impose the Higgs, stau, chargino mass bounds: $m_h > 113.5$ GeV, $m_{\tilde{\tau}} > 89$ GeV, $m_{\tilde{\chi}_1^\pm} > 103$ GeV, and also the 2σ constraint on the $b \rightarrow s\gamma$ branching ratio: $\text{Br}(B \rightarrow X_s\gamma) = (2.2-4.1) \times 10^{-4}$. It has been noted that the AMSB model is severely constrained by the recent measurement of the muon anomalous magnetic moment a_μ once we require that the conventional one-loop SUSY contribution $a_\mu^{\text{SUSY}} \gtrsim 10^{-9}$ which was taken as the 2σ lower bound [13]. Here we do *not* take this as a real constraint since the uncertainty in the hadronic contributions to a_μ can be as large as 10^{-9} [14]. Although not taken as a constraint, we specify the parameter region with $a_\mu^{\text{SUSY}} \gtrsim 10^{-9}$ or 5×10^{-10} for the completeness.

To conclude, we have examined the possibility of an observable same sign dilepton signal induced by the $\tilde{\nu}$ - $\tilde{\nu}^*$ mixing in the AMSB model. It is pointed out that the AMSB model has a good potential to provide an observable rate of signals since the mixing rate is naturally enhanced by $m_{3/2}/m_{\tilde{\nu}} = \mathcal{O}(4\pi/\alpha)$ while the sneutrino decay rate is small enough on the sizable portion of the phenomenologically allowed parameter space of the model. Our results depicted in Figs. 1 and 2 show that this is indeed the case for the neutrino masses suggested by the atmospheric and solar neutrino data. It is noted also that the same-sign dilepton signals can be used to determine $(m_\nu)_{ee}$ and $(m_\nu)_{\mu\mu}$, providing useful information on the neutrino masses and mixing angles which are complementary to those from neutrino experiments.

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REFERENCES

- [1] See for instance M. C. Gonzalez-Garcia et al., hep-ph/0009350; K. Choi et al., hep-ph/0107083.
- [2] M. Hirsch, H. V. Klapdor-Kleingrothaus and S. G. Kovalenko, Phys. Lett. B **398**, 311 (1997).
- [3] Y. Grossman and H. Haber, Phys. Rev. Lett. **78**, 3438 (1997).
- [4] Y. Grossman and H. Haber, Phys. Rev. D **59**, 093008 (1999); *ibid.* **63**, 075011 (2001).
- [5] E. J. Chun, hep-ph/0105157.
- [6] L. Randall and R. Sundrum, Nucl. Phys. **B557**, 79 (1999); G. F. Giudice, M. A. Luty, H. Murayama, and R. Rattazzi, J. High Energy Phys. **12**, 027 (1998); A. Pomarol and R. Rattazzi, *ibid.* **05**, 013 (1999); J. A. Bagger, T. Moroi, and E. Poppitz, *ibid.* **04**, 009 (2000).
- [7] M. Gell-Mann, P. Ramond and R. Slansky, in *Supergravity*, North-Holland, Amsterdam (1979); T. Yanagida, KEK Report No. 79-18 (1979).
- [8] H. P. Nilles, Phys. Rep. **150**, 1 (1984).
- [9] G. F. Giudice and R. Rattazzi, Phys. Rep. **322**, 419 (1999), and references therein.
- [10] T. Gherghetta, G. F. Giudice, and J. D. Wells, Nucl. Phys. **559**, 27 (1999); J. L. Feng and T. Moroi, Phys. Rev. D **61**, 095004 (2000).
- [11] K. Choi and A. Santamaria, Phys. Lett. B **267**, 504 (1991); E. Ma and U. Sarkar, Phys. Rev. Lett. **80**, 5716 (1998).
- [12] K. Choi, K. Hwang and W. Y. Song, preprint in preparation.
- [13] J. L. Feng and K. T. Matchev, Phys. Rev. Lett. **86**, 3480 (2001); K. Choi, K. Hwang, S. K. Kang, K. Y. Lee, and W. Y. Song, Phys. Rev. D **64**, 055001 (2001).

FIGURES

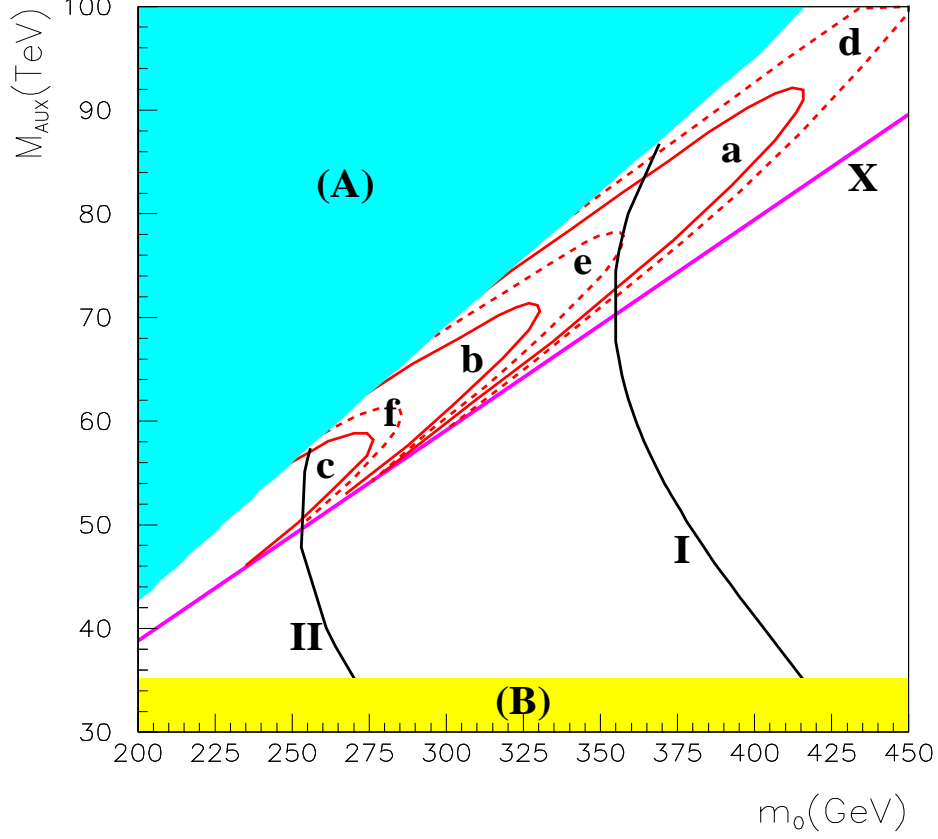


FIG. 1. Parameter regions with $\tan\beta = 5$ yielding $N_\mu \geq 20$ [inside the contour a], 10^2 [b], 5×10^2 [c] for $(m_\nu)_{\mu\mu} = 3 \times 10^{-2}$ eV; $N_\mu \geq 10^2$ [d], 5×10^2 [e], 2×10^3 [f] for $(m_\nu)_{\mu\mu} = 0.3$ eV. (A) and (B) represent the parameter regions forbidden by the stau and chargino mass bounds, respectively. Upper side of the line X denotes the region of LSP stau. Left sides of the lines (I) and (II) correspond to the region with $a_\mu^{\text{SUSY}} \geq 5 \times 10^{-10}$ and 10^{-9} , respectively.

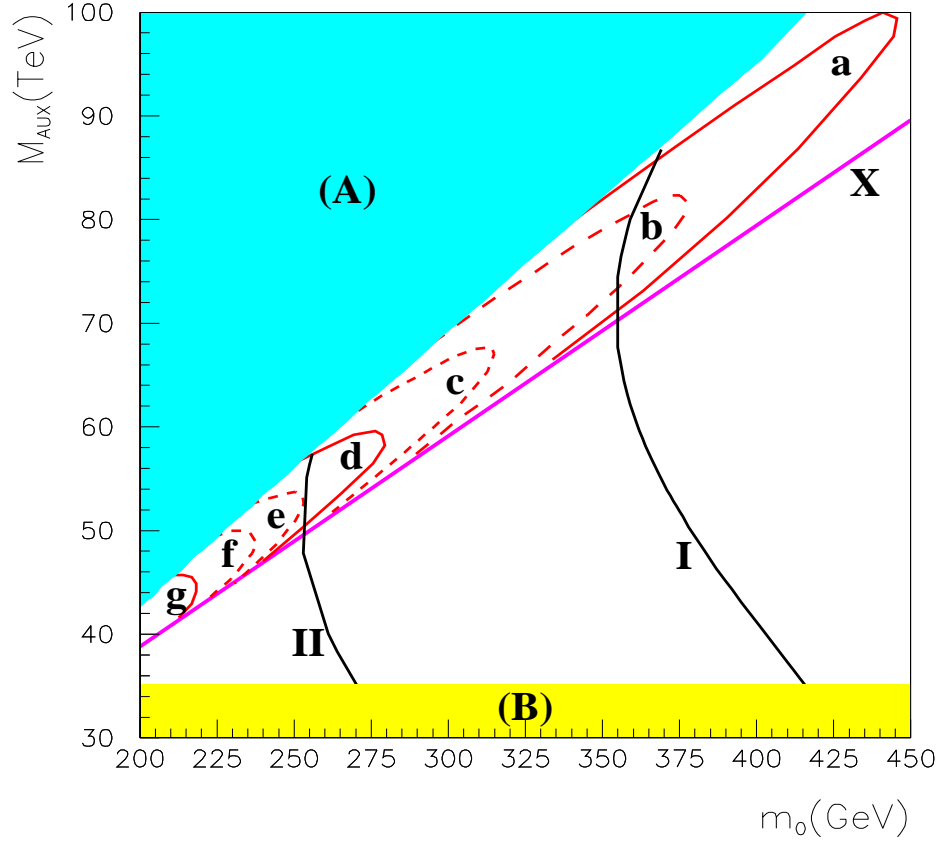


FIG. 2. Parameter regions with $\tan\beta = 5$ yielding $N_e \geq 20$ [a], 10^2 [b], 5×10^2 [c] for $(m_\nu)_{ee} = 10^{-2}$ eV; $N_e \geq 20$ [d], 10^2 [e], 5×10^2 [f] for $(m_\nu)_{ee} = 10^{-3}$ eV; $N_e \geq 10^2$ [g] for $(m_\nu)_{ee} = 10^{-4}$ eV.